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Baryogenesis from Domain Walls in the Next-to-Minimal Supersymmetric Standard Model

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Abstract

We consider the production of baryon number from collapsing domain walls, and in particular examine the magnitude of CP violation which is required in such schemes. Taking the conventional solution to the domain wall problem in the Next-to-Minimal Supersymmetric Standard Model as an example, we show that the observed baryon assymmetry of the universe may have been generated, even if the initial *explicit* CP violation in the Lagrangian were so small (*i.e.* gravitational) that it could never be experimentally detected. This is possible by having the *explicit* CP violation affect the way in which the walls collapse, rather than be responsible for the generation of baryon number directly. Net baryon number is created at the domain walls by the spontaneous breaking of CP.

1 Introduction

In models of baryogenesis at or close to the electroweak phase transition, one of the three Sakharov conditions [1] (no thermal equilibrium) is fulfilled by the phase transition itself, while the remaining two conditions (B, C and CP violation) are provided by sphaleron mediated processes and some extension to the Higgs sector usually involving an explicit CP violating phase [2, 3]. So far, the study of baryon number production involving topological defects has mainly addressed the first of these conditions. That is, it has been shown that the departure from thermal equilibrium may be provided by the collapse of cosmic strings or domain walls [4, 5]. In this paper we shall address the second. It is generally assumed that there must be an explicit CP violating extension to the Higgs sector similar in magnitude to that required for electro-weak baryogenesis, in order to bias the production of baryons. Such CP violating terms may eventually be detected through their contribution to the neutron electric dipole moment for example [6]. Here we shall show that this is not a necessary condition for baryogenesis from topological defects. In fact, provided that there is *spontaneous* breaking of CP when the domain walls form [7], it is possible to generate the observed baryon asymmetry with additional *explicit* CP violating terms which are gravitationally suppressed, and which will therefore never be detected. We stress the difference between 'spontaneous' violation of CP which is responsible for the local production baryon number at the domain walls, and the 'explicit' CP violation which is needed in order to have a global excess of baryons. In this respect, our picture is similar to that proposed in Ref.[8].

Our argument can be summarised as follows. Because the domain walls in question result from a breaking of CP, any particular wall is not CP invariant. Global CP invariance is provided by the fact that there exist different types of wall which are CP conjugates of each other. The mechanism which is invoked in order to remove the walls need not involve large terms in the Higgs potential. In fact for walls in which the higgs VEV, v, is of the order the electroweak scale (or larger), if the degeneracy in the minima is broken by gravitational couplings of order v^5/M_{pl}, the walls will disappear well before the onset of nucleosynthesis [9]. Since individual walls are not CP invariant, it is possible to generate a sufficient baryon number with explicit CP violation of order v^5/M_{pl}.

We shall demonstrate this using the next-to-minimal supersymmetric standard model (NMSSM). It should be borne in mind however that our discussion applies to any model in which the spontaneous breaking of CP produces domain walls. By choosing an example with three phases we are perhaps making things more difficult for ourselves, since models with more than two phases have their own special problems (some of which will be addressed in Ref.[10]). However it is interesting that a case with the required properties exists already in the literature.

The NMSSM [11] is an extension of the usual minimal supersymmetric standard model (MSSM) [12], in which the usual two Higgs doublets H_1 and H_2 which are necessary to give masses to the up and down type quarks are supplemented by a singlet Higgs superfield N. The usual μ term in the lagrangian, $\mu H_1 H_2$, is then eliminated by invoking a Z_3 symmetry under which every chiral superfield Φ transforms as $\Phi \rightarrow e^{2\pi i/3}\Phi$. The allowed terms in the superpotential are then $\lambda N H_1 H_2 - \frac{k}{3}N^3$, in addition to the usual fermion mass generating Yukawa terms, while the Higgs part of the soft supersymmetry breaking potential is extended by the inclusion of two more extra trilinear soft terms A_λ and A_k in place of the MSSM term $B\mu H_1 H_2$ to become

$$V_{soft}^{Higgs} = -\lambda A_\lambda (N H_1 H_2 + h.c.) - \frac{k}{3} A_k (N^3 + h.c.)$$

$$+m_{H_1}^2|H_1|^2 + m_{H_2}^2|H_2|^2 + m_N^2|N|^2 \quad (1)$$

where $H_1H_2 = H_1^0 H_2^0 - H^- H^+$, and we shall hereafter drop the 0 index for neutral Higgses.

The primary motivations for the NMSSM are the elimination (or at least reparametrisation) of the μ problem, which is that it is not clear what could be the origin of a μ parameter in $\mu H_1 H_2$ which is of order the electroweak scale; that it allows the evasion of the usual MSSM Higgs mass bounds [13]; and the fact that this relatively minor alteration to the model gives an extremely rich and complex Higgs and neutralino phenomenology which can be significantly different from that in the MSSM [14].

When electroweak symmetry breaking occurs, the three neutral CP-even Higgs scalars acquire VEVs. Using the parameters

$$\begin{aligned} H_1 &= \rho_1 e^{i\theta_1} \\ H_2 &= \rho_2 e^{i\theta_2} \\ N &= \rho_x e^{i\theta_x} \end{aligned} \quad (2)$$

it can be shown that any true minimum of the potential does not violate CP in the sense that the VEVs can always be made real by an appropriate field redefinition, up to the existence of three degenerate vacua related to each other by Z_3 transformations [15], and hence we have minima with $\theta_1 = \theta_2 = \theta_x = \frac{2\pi n i}{3}$ for integer n, [1] and with $\rho_i = \nu_i$, $\rho_x = x$. Note that although we have imposed that the ν_i be real, one or two of them may still be negative. We shall refer to the three minima

$$\begin{aligned} \theta_1 + \theta_2 &= 0, 2\pi/3, 4\pi/3 \\ \theta_x &= 0, 4\pi/3, 2\pi/3, \end{aligned} \quad (3)$$

as A, B, C respectively, and for convenience will assume that the evolution of the universe will ultimately end with phase A dominating.

After the electroweak phase transition the universe will be divided up into regions of different minima separated by domain walls. In each of the three degenerate minima, there is an operation which performs a CP transformation in the effective low energy theorem, and which maps the vacuum into itself. If we define the Z_3 operation to be $Z_3 : A \to B$, $Z_3 : B \to C$, and $Z_3 : C \to A$, then the transformations are

$$\begin{aligned} \mathrm{CP}_A &= \mathrm{CP} \\ \mathrm{CP}_B &= \mathrm{CP} Z_3 \\ \mathrm{CP}_C &= \mathrm{CP} Z_3^2, \end{aligned} \quad (4)$$

where here, CP is the transformation in the full theory. In each minimum, the two false vacua are CP conjugates of each other. Alternatively, we could have performed a field redefinition such that the true minimum (here A) is CP invariant.

2 Domain Walls in the NMSSM

Domain walls are one of the simplest types of topological defects [9], and form whenever the theory in question has a discrete number of degenerate vacua, usually due to the spontaneous

[1] We can use weak hypercharge to select any phase for $\theta_1 - \theta_2$.

breaking of a discrete symmetry. The simplest example occurs in the case of a real scalar field ϕ with potential $V = \lambda(\phi^2 - \nu^2)^2$. This potential clearly has two degenerate minima with $\phi = \pm\nu$ as a result of the Z_2 symmetry $\phi \to -\phi$. If we look for time independent solutions of the field equations which are translation invariant in two of the three space dimensions, and which obey $\phi(z = -\infty) = -\nu$ and $\phi(z = \infty) = \nu$ we find a solution $\phi = \tanh(z/\delta z)$. This is a domain wall, whose thickness δz is given by $\delta z = (\sqrt{2\lambda}\nu)^{-1}$, and which has a surface energy σ given by $3\sigma = 4\sqrt{2\lambda}\nu^3$. In more complicated models, it is no longer possible to solve the field equations analytically, but it is straightforward to solve them numerically. We find that $\delta z \sim \nu^{-1}$ and $\sigma \sim \nu^3$ as before, where ν is now some typical VEV of one of the fields, and the structure is in general similar.

Turning specifically to the NMSSM, the potential for the neutral scalars takes the form (at tree-level)

$$\begin{aligned} V = \frac{(g_1^2+g_2^2)}{8}(|H_1|^2 - |H_2|^2)^2 + \lambda^2|N|^2(|H_1|^2 + |H_2|^2) \\ +\lambda^2|H_1|^2|H_2|^2 - \lambda k(\bar{N}^2 H_1 H_2 + N^2 \bar{H}_1 \bar{H}_2) \\ +k^2|N|^4 + m_{H_1}^2|H_1|^2 + m_{H_2}^2|H_2|^2 + m_N^2|N|^2 \\ -\lambda A_\lambda(N H_1 H_2 + \bar{N}\bar{H}_1\bar{H}_2) - \frac{k}{3}A_k(N^3 + \bar{N}^3) \end{aligned} \tag{5}$$

With real VEVs $< \rho_1 >= \nu_1$, $< \rho_2 >= \nu_2$, $< \rho_x >= x$ our inputs are then $\tan\beta = \frac{\nu_2}{\nu_1}$, $r = \frac{x}{\nu}$, λ, k, A_λ, A_k, while $\nu^2 = \nu_1^2 + \nu_2^2$ is derived from the requirement that we have the correct Z mass. We choose to specify the VEVs as input parameters rather than the masses appearing in the potential for convenience, since we may immediately calculate the soft masses $m_{H_1}^2$, $m_{H_2}^2$, m_N^2 from the VEVs by using the minimisation conditions. Of course, this model typically has several different minima, usually including for example minima with only one of the three VEVs non-zero, and in order to study the vacuum structure it is necessary to find all of them to ensure that the minimum which we are analysing is indeed the deepest one.

Let us now turn to domain wall solutions of the field equations. These reduce to the six equations

$$\frac{d\phi_i}{dz} + \frac{1}{2}\frac{\partial V}{\partial \phi_i} = 0 \tag{6}$$

where ϕ_i is the real or imaginary part of one or other of the three scalar Higgs fields. We may then impose the boundary conditions that (H_1, H_2, N) are (ν_1, ν_2, x) at $z = -\infty$ and $(\nu_1 e^{2\pi i/3}, \nu_2 e^{2\pi i/3}, x e^{2\pi i/3})$ at $z = \infty$. It is straightforward to find solutions to such equations numerically, and by appropriate Z_3 field redefinitions it is clear that walls with the same structure exist between any two pairs of vacua.

A typical solution is shown in Figure 1, where we show the absolute values, phases, and energy density as a function of z in the wall region. The input parameters are $\lambda = k = 0.2$, $A_\lambda = A_k = 100$GeV, $\tan\beta = r = 2$. Here the total surface energy density of the wall is $7.1 \times 10^6\text{GeV}^3$ after we have subtracted the vacuum energy density, while the wall thickness is around 0.02GeV^{-1}, in reasonable agreement with the approximate arguments given above. It should be noted that even in the centre of the wall the VEVs are not zero, and so electroweak symmetry is not restored.

In fact, as the parameters are varied a very wide range of different behaviours and structures for the wall can be seen, with typically (for $\tan\beta > 1$ and $r > 1$) ρ_x remaining large over much of the region, while $\rho_2 > 0$ always but may become quite small near the centre of the

wall. The phase behaviour shown in Figure 1b, where the $U(1)_Y$ phase $\theta_1 - \theta_2$ goes from 2π to 0 continuously across the range is not universal but is typical. Unlike that for all the other variables, most of this change in $\theta_1 - \theta_2$ is outside the wall region, but we have explicitly checked that changing the size of the box does not have any significant impact on the total energy or the shape of the field configuration.

An example of a set of parameters for which electroweak symmetry is virtually restored is shown in Figure 2. Here $\lambda = k = 0.1$, $A_\lambda = A_k = 250\mathrm{GeV}$, $\tan\beta = 2$, $r = 5$. The total wall energy is $2.2 \times 10^7\mathrm{GeV}^3$, rather higher than before because the singlet VEV is larger, while the wall is now slightly wider. Although only ρ_1 is ever zero inside the wall, ρ_2 falls to under 3GeV, and is less than 10GeV for a region of width $\sim 0.04\mathrm{GeV}^{-1}$.

Of course, these are just two of a multitude of possible sets of parameters, each of which will give a wall with possibly very different characteristics. We also remark that there can be more than one solution for a given set of parameters, although for those shown there are no other wall solutions with higher surface energy. Since the VEVs of the higgs fields do not go through the origin, there is the possibility for 'double' (and also triple in this case) wall systems with the same phase on either side to form in the manner described in Ref.[16]. We shall assume that this does not occur, or at least that if it does the double walls are unstable to the formation of holes by quantum tunneling, which then expand under the surface tension destroying the wall. Our primary conclusion must be that for at least some sets of parameters, the domain walls possess exactly the properties which we will require in order to have them driving baryogenesis.

3 Baryogenesis from Z_3 Wall Networks

Having established this fact, let us go on to examine the possibilities for baryogenesis. After the phase transition we have an 'emulsion' of three phases separated by highly convoluted domain walls. CP is also spontaneously broken by the phase transition, but as yet there is no *explicit* CP violation. In fact, as we have seen, when going through a domain wall from $A \to C$, $B \to A$ or $C \to B$, the phase changes of the Higgs fields are equal and opposite to those occurring when going from $A \to B$, $B \to C$ or $C \to A$. We shall refer to these two types of transition as 'positive' and 'negative' respectively. The walls are not invariant under CP, and inside them the electroweak symmetry is restored if the vacuum expectation values of the Higgs fields vanishes (which, as we have seen, may or may not be the case depending on the details of the Higgs sector). Thus we shall assume that baryon number violating transitions will be in equilibrium in these regions, at plasma temperatures close to the phase transition. As domain walls move through space, the time-dependent change of phase of the Higgs fields occurring inside the walls will give rise to a non-zero chemical potential for baryon number and baryons will therefore be created.

Cosmology dictates that there is some mechanism which removes the walls and one suggestion, originally by Zel'dovich *et al* [9], is that the degeneracy of the vacua may be slightly broken, eventually leading to the dominance of the true vacuum. This point of view was recently supported by Rai and Senjanovic [17], who argue that gravitational interactions may explicitly break the discrete symmetries causing a slight non-degenaracy in the minima of the Higgs of order $\epsilon \sim v^5/M_{pl}$ (where v is a generic Higgs VEV of order M_W in this example). This suggestion was applied to this particular model in the context of string theories by Ellis *et al* [18].

We should point out two possible problems with this solution to the domain wall problem for this particular model. The first is the problem of destablising divergences which may generate

a large VEV for the singlet, and so destroy the solution which supersymmetry provides for the hierachy problem [19]. This is a potential fault in any model which includes gauge singlets. It is not clear which operators may be generated at the Planck scale or with what coefficients, but we note that the CP-violating gravitationally-suppressed operators which are necessary to remove domain walls do not in themselves generate such large singlet VEVs. Connected with this problem is the fact that if we break the Z_3 symmetry even by gravitational terms, we reintroduce the μ-problem since without the Z_3 symmetry there is nothing to prevent μ becoming large. These points detract from the NMSSM but are unavoidable; unless we are prepared to complicate the model by invoking inflation with reheating to a temperature less than the weak scale (and probably the Affleck-Dine mechanism for baryogenesis), we must certainly break the Z_3 symmetry explicitly. These are problems for the NMSSM as a whole and are secondary to our present more general aim of showing that domain walls can induce baryogenesis with small CP violation. We will not discuss them further, but will simply bear in mind that a full resolution of these problems of the NMSSM seems to require a greater understanding of the structure at the Planck scale.

The removal of the false vacua (and therefore the domain walls) proceeds as follows. For friction-free motion, the typical curvature scale, R, of the wall structure evolves roughly as the time for models with Z_N symmetry. Since we are not interested in the precise power law behaviour of the curvature scale, we shall neglect the conformal stretching due to the expansion of the universe. For detailed discussions of these points see Refs.[20, 21]. We shall also neglect the effects of friction on the motion of the walls. In fact this may be important at lower temperatures for domain walls associated with higgs fields. This is due to the walls' interaction with particles in the plasma, most importantly the bottom quarks, which are reflected off them with probability proportional to m_b^2/p^2 where p is the particle's momentum. Thus friction is unimportant for temperatures between E_W and 10 GeV. When the motion of the walls is friction dominated, they reach a terminal velocity determined by their curvature and by the density of the plasma. The typical curvature scale of the walls then increases as $t^{1/2}$ rather than t [22]. These points will be discussed in detail in [10].

Once the curvature scale has exceeded a critical value, *i.e.* when the pressure dominates over the tension, $\epsilon > \sigma/R$ where σ is the surface energy density, the domains of true vacuum begin to dominate and expand into the two domains of false vacuum. However, since CP is only broken spontaneously, any mechanism which removes the walls generates as much matter as anti-matter. In this case, the true vacuum, A, invades an equal area of B and C when it finally dominates (B and C must be degenerate if CP is not explicitly broken), and the production of baryons from negative walls exactly cancels that from positive ones.

Spontaneous CP violation *per se* is therefore not enough to generate a net baryon number. What is also required is some additional *explicit* CP violation in the Lagrangian, and this is where this paper differs from previous discussions. Previously attention has almost always been focussed on the biasing of the baryon number production directly at the collapsing domain walls (an exception being the scenario examined in Ref.[8], which bears some resemblances to this picture). Thus any explicit CP violation that was added to the Lagrangian was incorporated linearly into the production of baryon number. The resultant models required relatively large CP violating phases in the higgs sector.

However, even tiny (of order v^5/M_{pl}) CP violating terms will clearly effect the way the domain walls collapse, and, as argued in Ref.[17], there is no reason why gravitational terms that break the Z_3 should not also break CP. What we propose therefore, is that no two of the vacua are degenerate, so that C has a higher vacuum energy than B, which has a higher vacuum

energy than A. As the walls collapse therefore, A domains will invade both B and C, B domains will invade C but be invaded by A, and C will be invaded by both A and B. In order to show that this can generate a significant baryon asymmetry, consider the extreme case, in which the C phase has 'much' higher vacuum energy than B. Then the first pressure driven process to operate as the scale of the wall network increases is the collapse of C domains to be replaced by A and B. Both positive and negative walls will quickly accelerate to the speed of light, and the net baryon number generated will be close to zero. Now there will be only A and B phases left. The remaining B is finally removed when the non-degeneracy in A and B vacua becomes dominant. But the only walls which can do this are the positive $B \to A$ ones, and so there is clearly the potential for generating net baryons.

4 Simulation of the Z_3 Wall Networks

The number which we need in order to be able to estimate the baryon number is the average number of positive and negative walls which pass through a given point during the whole process. In order to show that this number can be close to one, we have simulated a Z_3 domain wall network evolving in 2 dimensions, in Minkowski space (*i.e.* neglecting the conformal stretching). We did this following Kawano [20]. First we begin with an arbitrary distribution of the three phases. The probability for each phase is $P_A = P_B = P_C = 0.33$ which interestingly is barely enough for them to percolate in three dimensions. Simulations on a cubic lattice give the percolation threshold to be 0.31 and simulations in continuum percolation theories give 0.295 ± 0.02 [23]. Thus the structure is expected to be tenuous and highly convoluted (*i.e.* 'spaghetti'–like). The walls are then divided into small lengths and released from the (in this case square) lattice. The evolution at each time step is determined by applying the equations of motion locally, taking the mass (proportional to the length) of the walls to be concentrated at the vertices between straight sections. In this we differ slightly from Kawano, who calculated the local curvature, since this enabled us to treat vertices with two and three walls attached on the same footing. In addition we did not include toroidal boundary conditions but let the ends of the walls slide along the edge of the box. We therefore do not expect our results to be accurate when the curvature scale is of the same order as the size of the box. Our basic unit for the simulation is shown in figure 3. A typical point, r_0, is connected to up to three other points. Each line has a perpendicular vector ϵ_{ij} associated with it which describes the magnitude and direction of the pressure acting on it.The rest mass of the vertex is given by half the sum of the lengths multiplied by the surface density σ;

$$m_0 = \frac{\sigma}{2} \sum_i |r_i - r_0|. \tag{7}$$

The force is given by $-\nabla E$ at the vertex

$$\frac{\partial p_0}{\partial s} = \sum_i \left(\gamma\sigma \frac{r_i - r_0}{|r_i - r_0|} + \epsilon_{0i} |r_i - r_0| \right), \tag{8}$$

where $\gamma = (1 - \dot{r}_0^2)^{-1/2}$ and s is proper time, so that the acceleration is given by

$$\frac{d^2 r_0}{dt^2} = \sum_i \left(2\gamma^{-2} \frac{r_i - r_0}{|r_i - r_0|} + \frac{\epsilon_{0i}\gamma^{-3}}{\sigma} |r_i - r_0| \right) / \sum_j |r_j - r_0|. \tag{9}$$

It is easy to verify that the continuous case is recovered in the limit as the size of the straight sections goes to zero. For example, consider the polygon made of N equal straight lengths, whose vertices are at r from the origin, and whose internal vacuum energy density is ϵ. The equation of motion above leads to

$$\frac{d^2r}{dt^2} = -\frac{(1-\dot{r}^2)}{r} - \frac{\epsilon\gamma^{-3}}{\sigma}\cos\frac{\pi}{N}, \tag{10}$$

which is that of a cylinder of radius r in the limit $N \to \infty$ [24]. It is convenient to scale everything in terms of the initial curvature scale R_0, so that $\rho_i = r_i/R_0$, and $\tau = t/R_0$, so that eq.(9) becomes,

$$\frac{d^2\rho_0}{dt^2} = \sum_i \left(2\gamma^{-2}\frac{\rho_i - \rho_0}{|\rho_i - \rho_0|} + \frac{\epsilon_{0i}R_0\gamma^{-3}}{\sigma}|\rho_i - \rho_0| \right) / \sum_j |\rho_j - \rho_0|. \tag{11}$$

The only free parameters in the simulation are therefore the two pressure variables, $\epsilon_B R_0/\sigma$ and $\epsilon_C R_0/\sigma$. All our results are presented with R_0 normalised to 1cm.

In figure 4 we can see how the network behaves without the effects of pressure. The scale of the structure evolves at roughly the speed of light (with about one wall per horizon) growing proportionally to the time. (This case, together with more general Z_N cases has been examined in more detail by Press *et al* [21].) The important point here is that (as remarked upon in Ref.[22]), without pressure, the evolution of the walls is mostly a question of topology. Those regions which are connected to two or four external walls tends to collapse, while those which are connected to eight or more external legs expand, due to the tension of the external legs pulling outwards. One can see this by considering any three leg vertex. A three leg vertex minimises its wall energy by trying to adopt a position in which the angles between the legs are equal and 120^o. Squares with four external legs do this by the vertices falling inwards trying to increase the internal 90^o angle. Octagons with eight external legs expand, trying to decrease a 135^o angle. Hexagonal structures (*i.e.* honeycombs) are stable. In fact for a general N sided polygon with N external legs, the equation of motion is easily found to be

$$\frac{\partial^2}{\partial\tau^2}r = -\frac{1}{r\gamma^2}\left(1 - \frac{1}{2\gamma\sin\left(\frac{\pi}{N}\right)}\right) + \frac{\epsilon R_0}{\sigma\gamma^3} \tag{12}$$

where we have normalised $r = R/R_0$, where R is the perpendicular distance from the centre to the edges of the polygon, and $\tau = t/R_0$. ϵ is the difference in vacuum energy between the inside and outside of the polygon, and R_0 is its initial size.

We now introduce pressure by switching on the ϵ above. This becomes dominant over the tension when

$$\left|\frac{\epsilon R}{\sigma}\right| \sim 1 \tag{13}$$

which for typical values of vacuum expectation values happens for $1cm < R_0 < 1m$ in cases where ϵ is induced gravitationally. The evolution of the system with pressure is shown in figure 5, where we have taken $\frac{\epsilon R_0}{\sigma} = 0, 0.25, 0.5$ for the phases A, B, C respectively. As mentioned earlier, the evolution of the network in terms of $r = R/R_0$ is the same for constant ϵR_0. As in the Z_2 case, larger structures are affected much more by the pressure. Since the structure of the walls is always increasing, once the pressure becomes dominant, collapse happens very quickly (between 10^{-10} and 10^{-8} seconds). Thus provided that the walls do not dominate in density

before they collapse, and also that the entropy released into the plasma is properly thermalised (both of which we shall assume), there is no danger of disturbing nucleosynthesis which begins at ~ 1 second.

The ratio of area cleared by positive walls to the total area we find to be 0.6, so that defining κ_{BG} as

$$\kappa_{BG} = \frac{\text{area of positive transitions - area of negative transitions}}{\text{total area}} \tag{14}$$

the global production of baryons is $\kappa_{BG} \approx 0.2$ of what it would be for maximal CP violation. For values of ϵ_c much larger than this, κ_{BG} rapidly approaches 1.

5 Discussion.

From this point, the analysis closely follows that in Ref.[4], and of course the same caveats apply. That is, we assume that the wall thickness is large enough to allow anomalous processes to occur. These may take the form of short range interactions of typical size $(g^2T)^{-1}$, where the electroweak symmetry is completely restored, or if the temperature is close to the phase transition one would expect sphaleron-like configurations straddling the domain wall to be possible. (Ideally one would like to be able to find these by constructing a non-contractible loop around the domain wall background.) Assuming that the sphaleron rate inside the walls is

$$\Gamma_B \sim \kappa(\alpha_W T)^4 \tag{15}$$

the final production of baryons is given by

$$\frac{n_B}{s} \approx \kappa\kappa_{BG} g_*^{-1} \alpha_W^4 \Delta(\theta_1 + \theta_2) \tag{16}$$

where g_* is the number of relativistic degrees of freedom contributing to the entropy, and $\Delta(\theta_1 + \theta_2)$ is the phase change of the higgs fields which is orthogonal to the Goldstone mode. This factor is to be multiplied by a coefficent of order 1 (see Ref.[3]). Note that in contrast to 'spontaneous baryogenesis' scenarios in which there is a possible suppression by a factor m_t^2/T^2 [25], here there is no suppression, because the relevant *equilibrium densities* to use are the ones obtaining outside the domain wall where m_t has its physical value. Notice that the production of baryon number is due to the translation of the walls, which contrasts with the case of electroweak strings, in which it is due to a decrease in the total volume covered by strings as they collapse; because of this there is no volume suppression factor (SF). The change in higgs phase is

$$\Delta(\theta_1 + \theta_2) \sim 2\pi/3, \tag{17}$$

so that

$$\frac{n_B}{s} \approx 10^{-8} \kappa\kappa_{BG} g_*^{-1}. \tag{18}$$

Bearing in mind our earlier discussion, there is the possibility of a much larger biasing of the potential. We note that this mechanism works for more general ϵ, provided firstly that the amount of explicit CP violation is of the same order as the explicit Z_3 breaking, and secondly that they are both not so large that the walls collapse immediately on forming. Thus we are in the novel position of being able to place (albeit extremely weak) lower and upper bounds on the amount of explicit CP violation allowed. The mechanism works only when the scale at which

the pressure dominates is larger than the size of the protodomains ($\approx (g_2^2 T)^{-1}$) during the phase transition which gives,

$$\epsilon \lesssim g_2^2 T_c \sigma. \tag{19}$$

For typical values of σ we find,

$$\epsilon \lesssim 10^7 \mathrm{GeV}^4 \tag{20}$$

which not surprisingly is just less than M_W^4. In addition we require that the temperature be close to the weak scale, for the anomalous processes to be operative inside the wall. However the domain sizes grow at speeds comparable to the speed of light. FRW cosmology gives t =2.42 $g_*^{-1/2}(T/\mathrm{MeV})^{-2}$ secs. For $T \sim M_W$ we find $t \sim 10^{-10}$ secs. Thus it is possible for 1 cm size structures to grow at temperatures close to the weak scale, implying that even CP violation induced by gravity could be the driving force behind baryon production for weak scale phase transitions.

In fact the following exercise is instructive. Suppose that the anomalous processes are effective down to a temperature $T_* < T_c$, and that the pressure and surface energy terms are given by

$$\epsilon = v^5/M \ ; \ \sigma = v^3 \tag{21}$$

where $v = \mathcal{O}(M_W)$ is the VEV of the higgs fields, and M is the mass scale of the physics which is responsible for the CP violation. Suppose also that the curvature scale increases at some sizeable fraction, β, of the speed of light, $R = \beta(t - t_c)$. Then in order for this mechanism to work, we require that the pressure is dominant over the surface tension for $t = t_*$,

$$\epsilon > \sigma/\beta(t_* - t_c). \tag{22}$$

This gives an upper bound on M,

$$M \lesssim M_{pl} \left(0.3\beta g_*^{-1/2} \left(\frac{v}{T_*}\right)^2 \left(1 - T_*^2/T_c^2\right)\right). \tag{23}$$

So unless T_* is extremely close to T_c, gravitational couplings could be responsible for this mechanism, giving a lower bound on ϵ of,

$$\epsilon \gtrsim 10^{-8} \mathrm{GeV}^4. \tag{24}$$

When this bound is saturated, on dimensional grounds one expects the contribution to the electric dipole moment of the explicit CP violation, to be of the order of $\delta d_n < 10^{-42} e$cm [6].

Finally, we note that this mechanism is possible for any model with a spontaneous CP breaking transition occuring at an energy scale, v, which is higher than the electroweak scale, provided that some domain walls remain at the time of the electroweak transition. In this case anomalous processes are guaranteed to be in equilibrium when the wall network collapses. The same considerations apply here. That is

$$g^2 v^4 \gtrsim \epsilon \gtrsim 10^{-8} \mathrm{GeV}^4. \tag{25}$$

In this case the lower bound is *less* than what would be expected to be induced by gravity, since we require simply that the domain walls collapse before the electroweak phase transition whilst anomalous processes are still in equilibrium.

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References

[1] A. D. Sakharov, Pisma Zh. Eksp. Teor. Fiz. **5** (1967) 32; (JETP Lett. 5 (1967) 24)

[2] A. G. Cohen and D. B.Kaplan, Phys. Lett. **B199** (1988) 251;
N. Turok and J. Zadrozny, Phys. Rev. Lett. **65** (1990) 1;
A. G. Cohen, D. B.Kaplan and A. E. Nelson, Phys. Lett. **B263** (1991) 86; Nuc. Phys. **B373** (1992) 453;
A. G. Cohen and A. E. Nelson, Phys. Lett. **B297** (1992);
G. F. Giudice and M. Shaposhnikov, Phys. Lett. **B326** (1993) 118

[3] S. A. Abel, W. N. Cottingham and I. B. Whittingham, Nuc. Phys. **B410** (1993) 173

[4] R. H. Brandenberger and A-C. Davis Phys. Lett. **B308** (1993) 79; Phys. Lett. **B332** (1994) 305;
R. H. Brandenberger, A-C. Davis and M. Trodden Phys. Lett. **B335** (1994) 123; hep-ph/9406355; hep-ph/9412266;
R. H. Brandenberger, A-C. Davis, T. Prokopec and M. Trodden hep-ph/9409281;
R. H. Brandenberger, A-C. Davis and M. J. Rees hep-ph/9501040

[5] H. Lew and A. Riotto, Phys. Lett. **B309** (1993) 258

[6] A. Kazarian, S. Kuzmin and M. Shaposhnikov, Phys. Lett. **B276** (1992) 131;
A. Cohen and A. Nelson, Phys. Lett. **B279** (1992) 111;
R. Mohapatra and X. Zhang, Phys. Rev. **D46** (1992) 5331;
X. Zhang and B.-L. Young, Phys. Rev. **D49** (1994) 563

[7] T. D. Lee, Phys. Rep. **9**, No. 2 (1974) 143

[8] D. Comelli, M. Pietroni and A. Riotto, Phys. Lett. **B306** (1993) 67; Nucl. Phys. **B412** (1994) 441; Phys. Rev. **D50** (1994) 7703; Phys. Lett. **B343** (1995) 207

[9] Ya. B. Zel'dovich, I. Yu. Kobzarev and L. B. Okun, Zh. Eksp. Teor. Fiz. **67** (1974) 3;
T. W. B. Kibble, J. Phys. **A9** (1976) 1387; Phys. Rep. **67** (1980) 183;
A. Vilenkin, Phys. Rep. **121** (1985) 263

[10] S. A. Abel, S. Sarkar and P. L. White, in preparation

[11] P. Fayet, Nucl. Phys. **B90** (1975) 104;
H.-P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. **B120** (1983) 346;
J.-M. Frere, D. R. T. Jones and S. Raby, Nucl. Phys. **B222** (1983) 11;
J.-P. Derendinger and C. A. Savoy, Nucl. Phys. **B237** (1984) 307;
L. Durand and J. L. Lopez, Phys. Lett. **B217** (1989) 463;
M. Drees, Intern. J. Mod. Phys. **A4** (1989) 3645;
J. Ellis, J. Gunion, H. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. **D39** (1989) 844

[12] For reviews see H.-P. Nilles, Phys. Rep. **110** (1984) 1;
H. E. Haber and G. L. Kane, Phys. Rep. **117** (1985) 75

[13] G. Kane, C. Kolda, J. Wells, Phys. Rev. Lett. **70** (1993) 2686;
J. Espinosa, M. Quiros, Phys. Lett. **B302** (1993) 51;

J. Espinosa, M. Quiros, Phys. Lett. **B279** (1992) 92;
U. Ellwanger and M. Rausch de Traubenberg, Z. Phys **C53** (1992) 521;
U. Ellwanger and M. Lindner, Phys. Lett. **B301** (1993) 365;
U. Ellwanger, Phys. Lett. **B303** (1993) 271;
P. N. Pandita, Phys. Lett. **B318** (1993) 338; Z. Phys. **C59** (1993) 575;
T. Elliott, S. F. King and P. L. White, Phys. Lett. **B305** (1993) 71; **B314** (1993) 56.

[14] J. F. Gunion, H. E. Haber, Nucl. Phys. **B278** (1986) 449;
T. Elliott, S. F. King and P. L. White, Phys. Rev. **D49** (1994) 2435;
U. Ellwanger, M. Rausch de Traubenberg and C. A. Savoy, LPTHE Orsay 95-04, LPT Strasbourg 95-01, SPhT Saclay T95/04;
F. Franke, H. Fraas, Phys. Lett. **B336** (1994) 415;
F. Franke, H. Fraas, Preprint WUE-ITP-95-003, hep-ph/9504279;
K. Olive, D. Thomas, Nucl. Phys. **B335** (1991) 192;
R. A. Flores, K. Olive, D. Thomas, Phys. Lett. **B263** (1991) 425;
S. A. Abel, S. Sarkar, I. B. Whittingham, Nucl. Phys. **B392** (1993) 83.

[15] J. C. Romão, Phys. Lett. **B173** (1986) 309

[16] G. Dvali, Z. Tavartkiladze and J. Nanobashvili, IFUP-TH 60/94, hep-ph/9411387

[17] B. Holdom, Phys. Rev. **D28** (1983) 1419; B. Rai and G. Senjanovic, Phys. Rev. **D49** (1994) 2729

[18] J. Ellis, K. Enqvist, D. V. Nanopoulos, K. Olive, M. Quiros and F. Zwirner, Phys. Lett. **B176** (1986) 403

[19] J. Bagger and E. Poppitz, Phys. Rev. Lett. **71** (1993) 2380

[20] L. Kawano, Phys. Rev. **D41** (1990) 1013

[21] W. H. Press, B. S. Ryden, D. N. Spergel, Ap. J. **347** (1989) 590; Ap. J. **357** (1990) 293

[22] I. M. Lifshitz, Sov. Phys. JETP **15** (1962) 939

[23] D. Stauffer, Phys. Rep. **54** (1979) 1; B. Lorenz, I. Orgzall and H.-O. Heuer, J. Phys. **A26** (1993) 4711

[24] L. Widrow, Phys. Rev. **D39** (1989) 3571; Phys. Rev. **D39** (1989) 3576; Phys. Rev. **D40** (1989) 1002

[25] M. Dine and S. Thomas Phys. Lett. **B328** (1994) 73

Figure Captions

Figure 1a

Absolute magnitures of the fields as a function of position for the parameters $\lambda = k = 0.2$, $A_\lambda = A_k = 100$GeV, $\tan\beta = 2$, $r = 2$. The three lines show, from top to bottom, ρ_3, ρ_2, ρ_1.

Figure 1b

Phases of the fields as a function of position for the same parameters as Figure 1a. The three lines show $\theta_+ = \theta_1 + \theta_2$ (solid lines) , $\theta_- = \theta_1 - \theta_2$ (long dashed lines), θ_x (short dashed lines).

Figure 1c

Surface energy density of the wall as a function of position for the same parameters as Figure 1a.

Figure 2a

Absolute magnitures of the fields as a function of position for the parameters $\lambda = k = 0.1$, $A_\lambda = A_k = 250$GeV, $\tan\beta = 2$, $r = 5$. The three lines show, from top to bottom, ρ_3, ρ_2, ρ_1.

Figure 2b

Phases of the fields as a function of position for the same parameters as Figure 2a. The three lines show $\theta_+ = \theta_1 + \theta_2$ (solid lines) , $\theta_- = \theta_1 - \theta_2$ (long dashed lines), θ_x (short dashed lines).

Figure 3

Basic unit for wall simulation.

Figure 4

Wall evolution without pressure. The four time-slices shown have time 0.5 10^{-10}s, 1.5 10^{-10}s, 2.5 10^{-10}s, 3.75 10^{-10}s, for upper left, upper right, lower left, lower right respectively. κ_{BG} is less than 0.01 always.

Figure 5

Wall evolution without pressure. The four time-slices shown have time 0.6 10^{-10}s, 1.5 10^{-10}s, 2.4 10^{-10}s, 3.75 10^{-10}s, for upper left, upper right, lower left, lower right respectively. κ_{BG} is -0.009, -0.004, 0.023, 0.104. Final value of κ_{BG} after all walls disappeared was around 0.2.

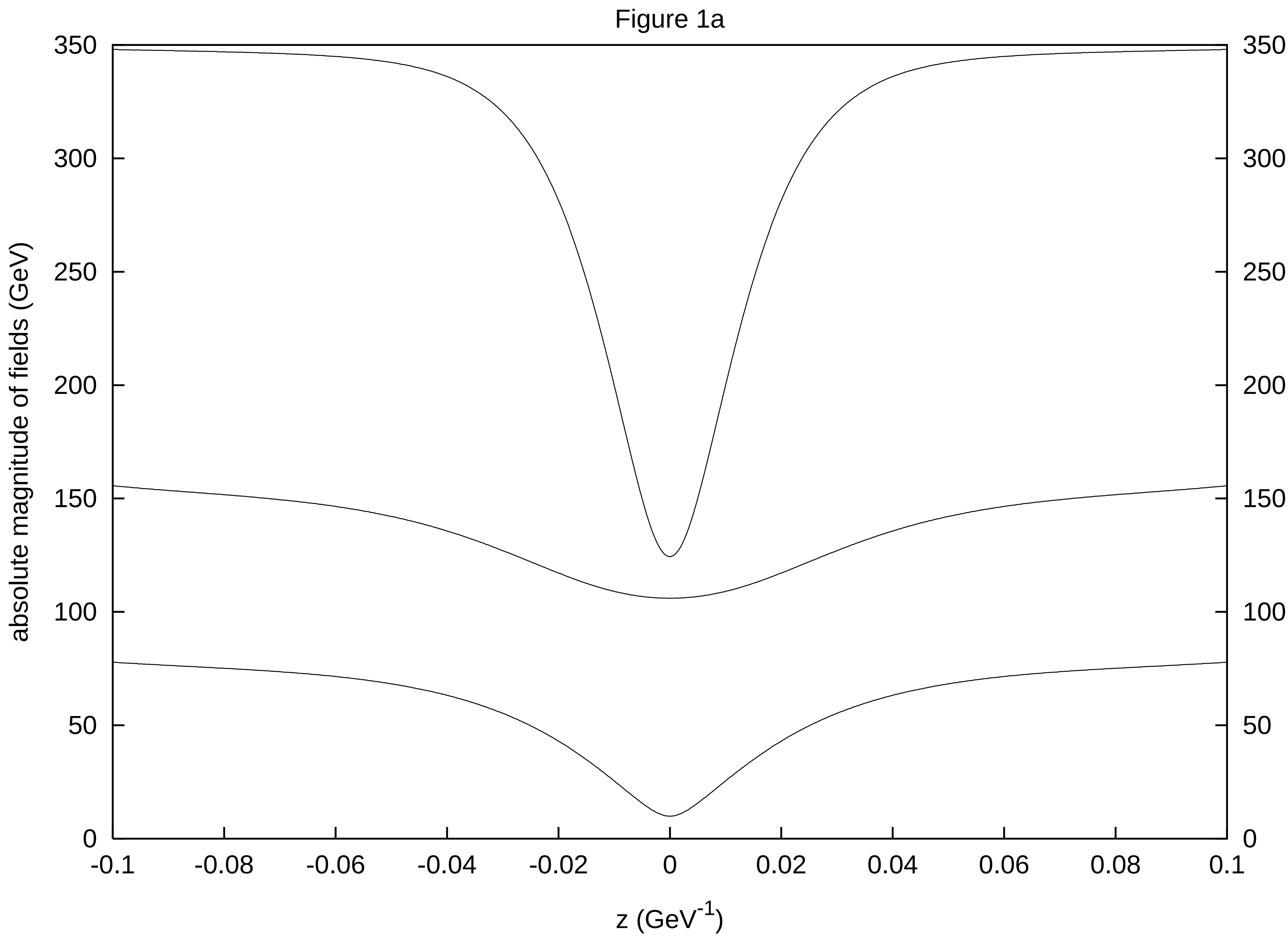
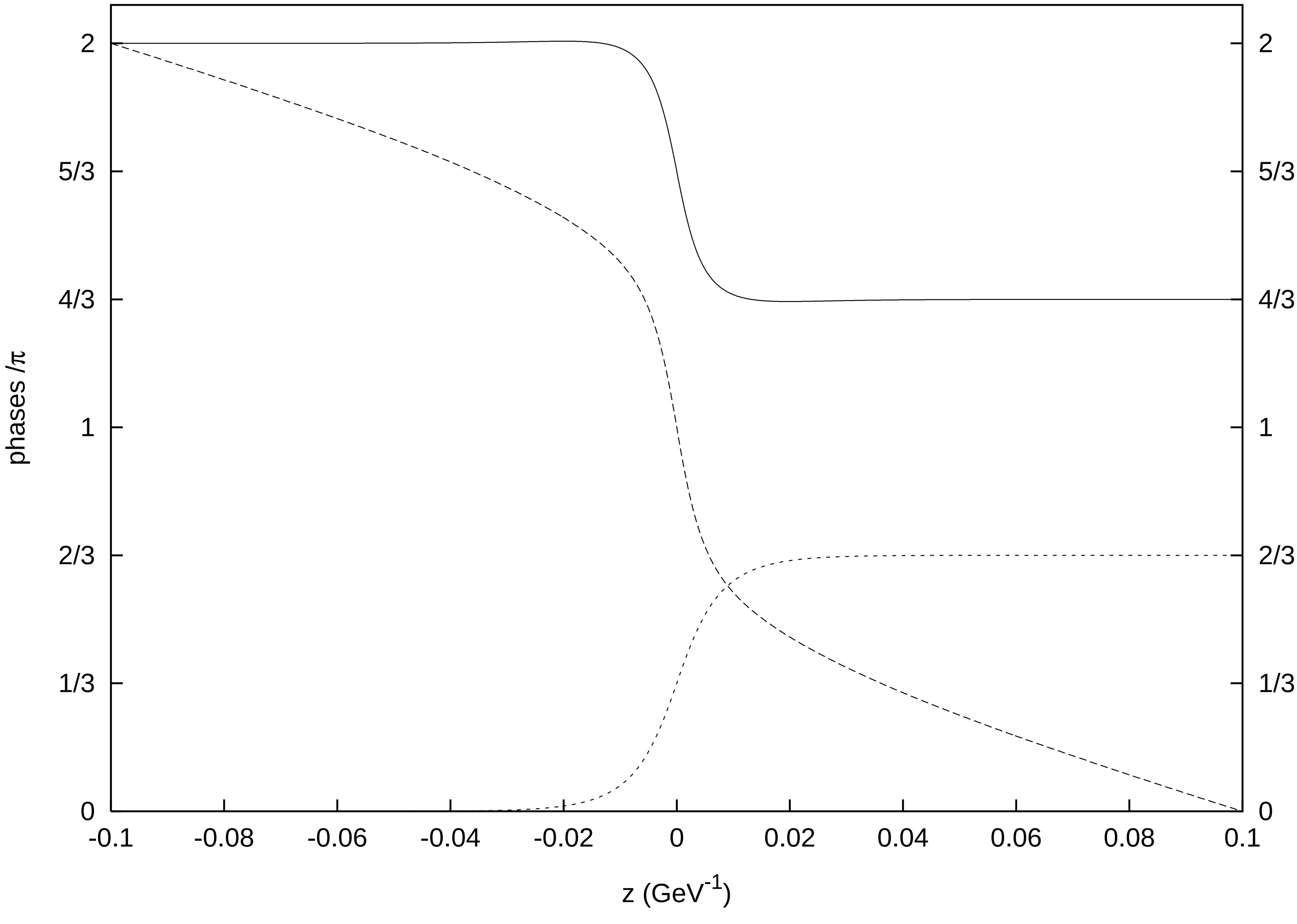


Figure 1a
absolute magnitude of fields (GeV)
z (GeV^{-1})
0
50
100
150
200
250
300
350
-0.1
-0.08
-0.06
-0.04
-0.02
0
0.02
0.04
0.06
0.08
0.1

Figure 1b



phases /π
2
5/3
4/3
1
2/3
1/3
0
-0.1
-0.08
-0.06
-0.04
-0.02
0
0.02
0.04
0.06
0.08
0.1
z (GeV^{-1})

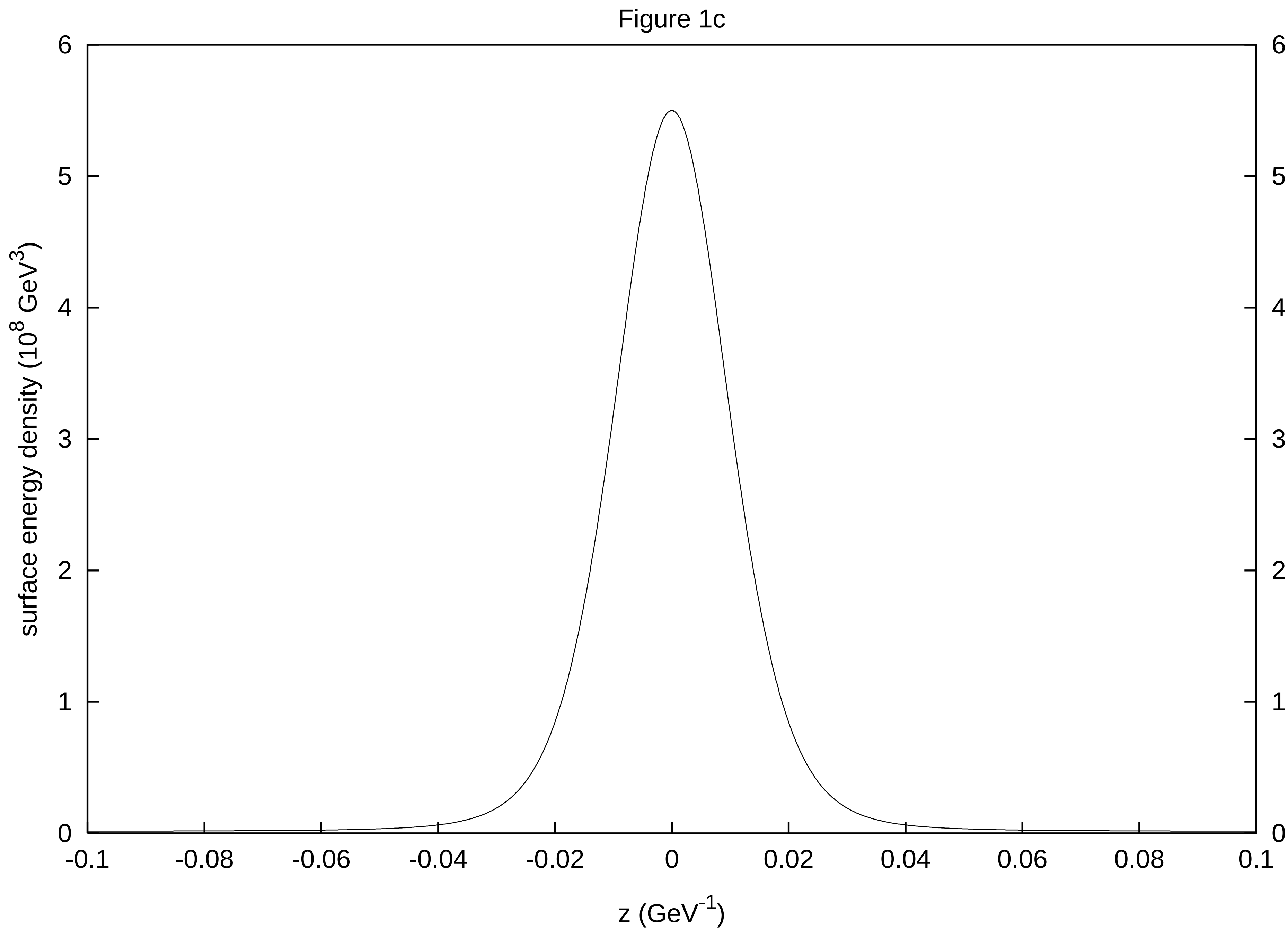


Figure 1c
surface energy density (10^8 GeV^3)
z (GeV^{-1})
0
1
2
3
4
5
6
-0.1
-0.08
-0.06
-0.04
-0.02
0
0.02
0.04
0.06
0.08
0.1

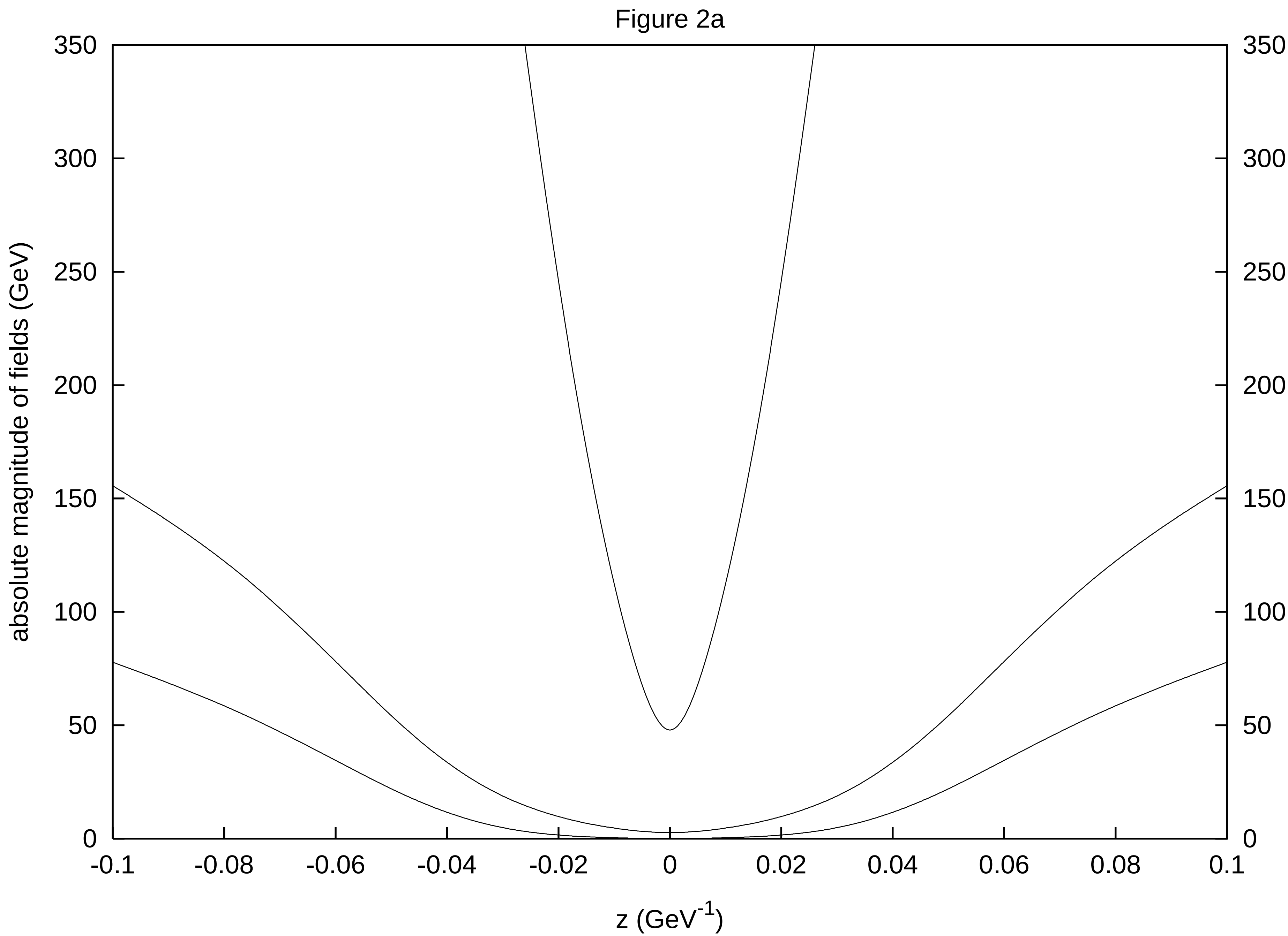
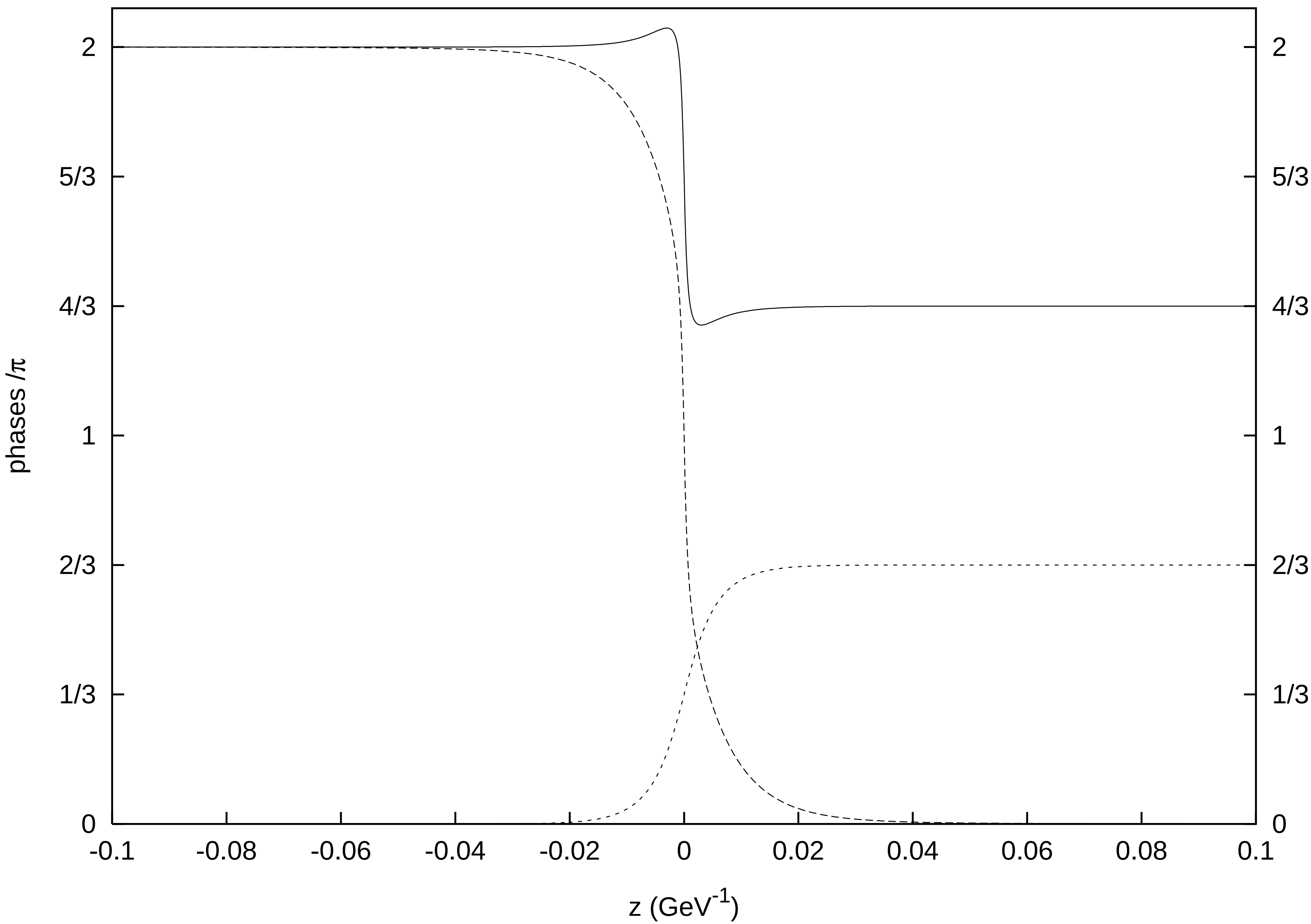


Figure 2a
absolute magnitude of fields (GeV)
350
300
250
200
150
100
50
0
-0.1
-0.08
-0.06
-0.04
-0.02
0
0.02
0.04
0.06
0.08
0.1
z (GeV^{-1})

Figure 2b



phases /π
z (GeV-1)
2
5/3
4/3
1
2/3
1/3
0
-0.1
-0.08
-0.06
-0.04
-0.02
0
0.02
0.04
0.06
0.08
0.1

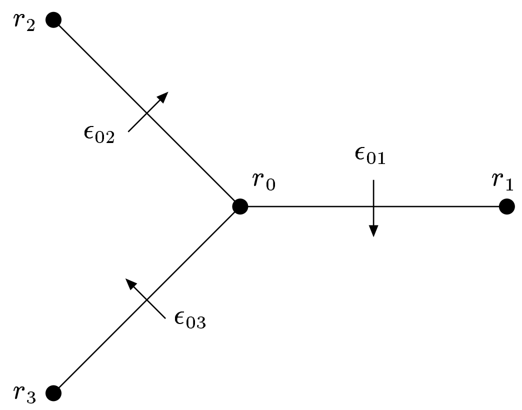


Figure 3

Figure 4

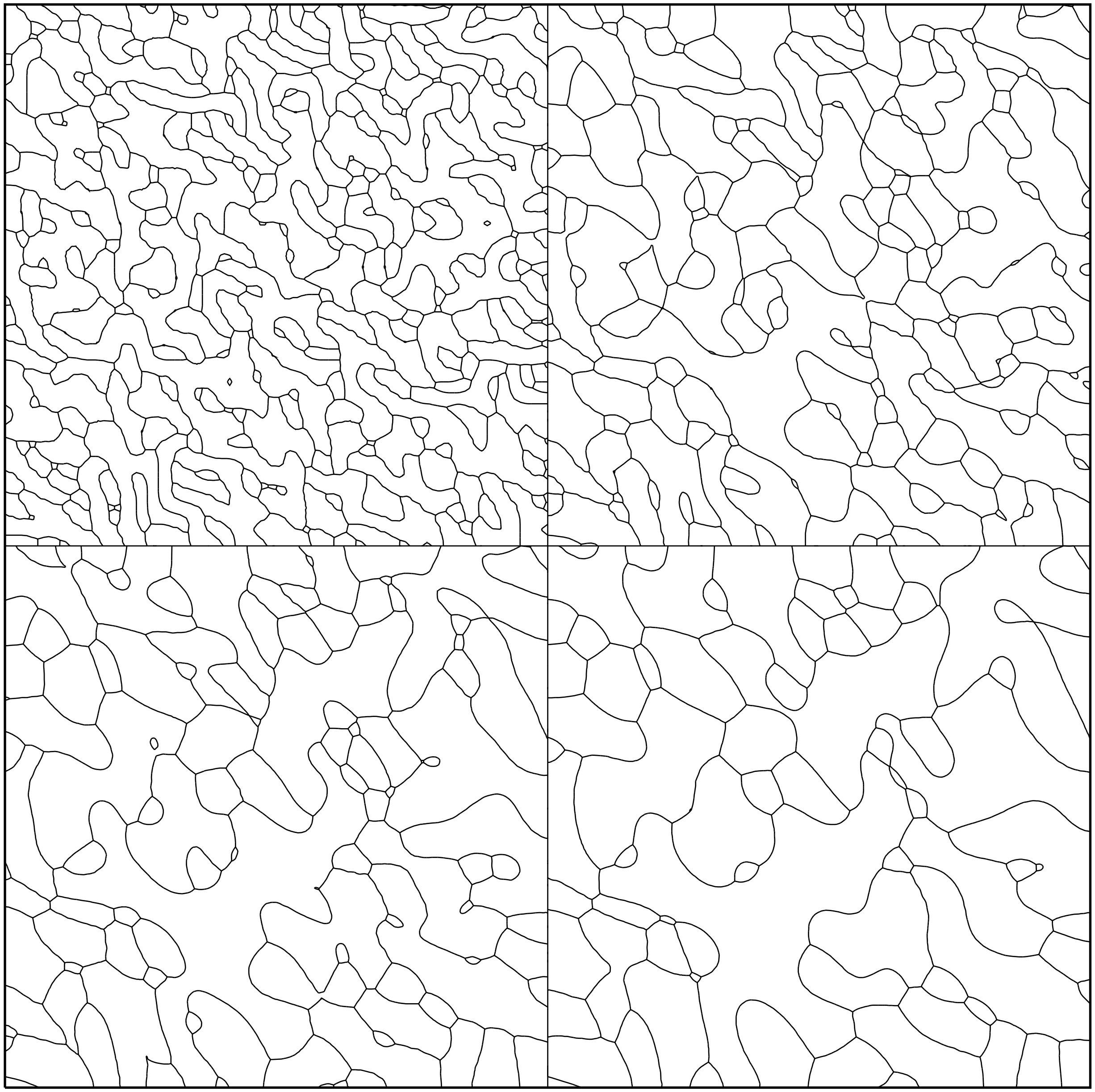


Figure 5